PHYSICAL REVIEW B VOLUME 42, NUMBER 1 1 JULY 1990

Tilt-modulus enhancement of the vortex lattice in the layered superconductor $2H$-$NbSe_2$

P. Koorevaar, J. Aarts, P. Berghuis, and P. H. Kes
Kamerlingh Onnes Laboratory, Leiden University, P.O. Box 9506, 2300 RA Leiden, The Netherlands
(Received 11 April 1990)

The field dependence of the pinning force has been studied in thin single crystals of the layered superconductor $2H$-$NbSe_2$ in fields directed perpendicular to the layers. At high fields a peak effect is observed which sets in at about $B_{co} \approx 0.8 B_{c2}$. Below this field the pinning force agrees well with the theory of two-dimensional collective pinning. The onset of the peak is triggered by the transition to three-dimensional flux-line lattice (FLL) disorder at the field B_{co}. Comparison of the crossover field with the criterion set by the collective-pinning theory reveals that the tilt modulus of the FLL in a layered superconductor is considerably reduced. The reduction factor corresponds very well to recent theoretical predictions. These results are of importance for the prediction of depinning and flux-line lattice melting in all kinds of anisotropic superconductors.

The discovery of the high-temperature superconductors (HTS) has greatly stimulated the study of flux-line lattice (FLL) properties. One of the many interesting predictions[1] is the dependence of the nonlocal tilt modulus of the FLL c_{44} on the anisotropy parameter Γ of layered superconductors. Here Γ is the ratio between the effective electron masses perpendicular (m_z) and parallel (m) to the layers, $\Gamma = m_z/m$. In essence, with increasing Γ, the flux lines become more disklike[2] leading to a reduction of c_{44}. FLL entanglement,[3] FLL melting,[1,4] "giant" flux creep,[5] depinning temperature,[6] and vortex glass-to-liquid phase transition[7,8] are among the new issues that are related to the peculiar characteristics of the extremely anisotropic HTS.

It is of great importance for the study of all these phenomena to test the predictions given by Houghton, Pelcovits, and Subdø.[1] Information on the actual value of c_{44} can be obtained from the field and temperature at which the crossover takes place from a two-dimensionally (2D) disordered (straight) FLL to a three-dimensionally disordered spaghetti state.[9,10] Such a transition reveals itself by a sudden increase of the pinning force F_p (peak effect) at the crossover field $B_{co} = b_{co} B_{c2}$. The most convenient experimental conditions for observing the crossover are obtained for weak-pinning thin-film superconductors, so that the behavior for $B < B_{co}$ is well described by the 2D collective-pinning (2DCP) theory[11,12] and $B_{co} \approx 0.8 B_{c2}$ is within experimental range. The HTS themselves are not well suited for these investigations because of strong flux-creep effects and therefore we concentrated our research on a low-temperature layered superconductor, $2H$-$NbSe_2$, which is an anisotropic type-II superconductor with $T_c \approx 7.3$ K. The structure has been well characterized[13] and consists of Nb planes sandwiched between two Se planes. The anisotropy of the superconducting parameters has been the subject of several studies.[14,15] The Ginzburg-Landau (GL) coherence length perpendicular to the layers is much larger than the distance between two superconducting Nb layers. Therefore the 3D GL theory for anisotropic superconductors describes the phenomena reasonably well.

For randomly distributed pinning centers (pins) no long-range order (LRO) is present in the FLL.[11] Short-range order remains in elastically independent correlated regions with volume $V_c = R_c^2 L_c$, where R_c and L_c are the transverse and longitudinal correlation lengths, respectively. The strength of the pinning is given by the parameter W which depends on the concentration of the pins and their elementary interaction with the FLL. The pinning-force density F_p and the critical-current density J_c follow from

$$F_p = B J_c = (W/V_c)^{1/2} . \qquad (1)$$

If elastic deformations are predominant the correlation lengths can be derived from the balance of elastic energy and pin energy. This situation only occurs if the disorder in the direction of the field can be neglected.[9,12] In that case $L_c = d$ (this is the 2DCP case) and R_c is given by

$$R_c = r_f c_{66} \{ 8\pi d / [W \ln(w/R_c)] \}^{1/2} . \qquad (2)$$

Here r_f is the range of the elementary pinning force, w the width, and d the thickness of a thin film with the field applied perpendicular to the film, and c_{66} is the shear modulus of the FLL, which according to Brandt, is given by

$$c_{66} = (B_c^2/4\mu_0) b (1 - 0.29b)(1-b)^2$$

for large κ (GL parameter) superconductors. B_c is the thermodynamic critical field and $b = B/B_{c2}$. For fields $b > 0.2$ one can put $r_f \approx a_0/2$. The situation of 2DCP is characterized by $F_p \propto B^{1/2}$ for low fields. For thin films a small peak effect is observed close to B_{c2} as soon as plastic deformations become predominant. For thick films a dimensional crossover (DCO) to 3D disorder results in a steep increase of F_p caused by a highly defective FLL.

It has been determined empirically that the DCO sets in when

$$L_c = (c_{44}/c_{66})^{1/2} R_c = d/2 , \qquad (3)$$

where c_{44} is the tilt modulus of the FLL. Brandt[16] demonstrated that c_{44} depends on the wave vector $\mathbf{k} = (k_\perp = (k_x^2 + k_y^2)^{1/2}, k_z)$ of the deformation field. For isotropic superconductors

$$c_{44}(k_\perp,k_z)\approx(B^2/\mu_0)k_h^2/(k_\perp^2+k_z^2+k_h^2)$$

with $k_h^2=0.86(1-b)/\lambda^2$ and λ the penetration depth. It has been explained in Ref. 17 that at the DCO $k_\perp \ll k_B$ and $k_z\approx\xi^{-1}$, so that $k_z \gg k_B$ and $k_z \gg k_h$. Here k_B is the radius of the Brillouin zone in the circular cell approximation $k_B^2=8\pi/a_0^2\sqrt{3}$ and $a_0=(2/\sqrt{3})^{1/2}(\phi_0/B)^{1/2}$ is the FLL parameter. The condition $k_z\approx\xi^{-1}$ reflects that the core-size of a screw dislocation is of the order ξ, and that, as soon as L_c related to the smallest possible wavelength fulfills the DCO criterion, screw dislocations are spontaneously created, probably at the surface. Substituting c_{44} and c_{66} in (3) yields for an isotropic superconductor that $(L_c/R_c)_{\rm co}\approx 3[b_{\rm co}/(1-b_{\rm co})]^{1/2}$ at the DCO.

For anisotropic superconductors c_{44} is much softer and is given by[1]

$$c_{44}(k_\perp,k_z)=c_{44}(0)\frac{m}{m_z}\times\left[\frac{k_h^2}{k_\perp^2+(m/m_z)(k_h^2+k_z^2)}+\frac{1-b}{2b\kappa^2}\right], \quad (4)$$

with $c_{44}(0)=B^2/\mu_0$. The ratio m/m_z is experimentally obtained from the upper-critical-field slopes in both orientations using $m/m_z=(S_\perp/S_\parallel)^2$ with $S_{\parallel,\perp}=-dB_{c2\parallel,\perp}/dT|_{T_c}$. As we will see below for the $NbSe_2$ samples under investigation, $R_c(b_{\rm co})\gg a_0$ at the DCO, so that the most important deformations have wave vectors $k_\perp^2 \ll k_B^2$. For k_z we should now substitute ξ_z^{-1}, where ξ_z is the GL coherence length perpendicular to the Nb layers. Substitution in (4) and (3) yields both c_{44} and L_c at the DCO,

$$c_{44,\rm co}=(B_c^2/\mu_0)(m/m_z)(1+1.72b_{\rm co})b_{\rm co}(1-b_{\rm co})$$

and

$$\frac{L_{c,\rm co}}{R_c(b_{\rm co})}=2\left(\frac{m}{m_z}\right)^{1/2}\left[\frac{1+1.72b_{\rm co}}{(1-0.29b_{\rm co})(1-b_{\rm co})}\right]^{1/2}\approx 3.93\left(\frac{m}{m_z}\right)^{1/2}\left(\frac{b_{\rm co}}{1-b_{\rm co}}\right)^{1/2}. \quad (5)$$

For $R_c(b_{\rm co})$ the value following from (2) should be used. The approximation of the prefactor 3.93 is better than 1% for $b>0.7$. The crossover field follows from $L_c=d/2$.

In order to enable the observation of 2DCP we repeatedly cleaved as-grown $2H$-$NbSe_2$ single crystals to obtain the smallest possible thickness. A well-known procedure was used of sandwiching the crystals between tape stuck onto two object slides. Taking the slides apart cleaves the crystals without causing significant surface damage as far as light-microscope inspection can show. The tape was removed by solving it in toluene for about 12 h. This provided us with crystals ranging in thickness between 2 and 50 μm. Bar-shaped samples were prepared by cutting the thinnest crystals to a typical size of 1×5 mm. The measurements discussed in this paper were carried out on a sample with $d=2$ μm (sample I) and on one with $d=15.5$ μm (sample II). Contacts for four-probe resistance measurements were made by silver paint.

The T_c's determined from the midpoints of the R vs T curves of samples I and II are 7.26 and 7.35 K, respectively. The transition widths ΔT_c of 150 mK were defined by the intercepts between the 10% and 90% normal-state resistances. Typical values of $T_c=7.18$–7.39 K reported by Toyota *et al.*[14] are close to ours. The residual resistance ratio of our samples is $R(T=293\text{ K})/R(T=7.5\text{ K})=34$, the room-temperature resistivity is 110×10^{-8} Ω m.

Upper critical fields were determined from R vs B transitions at constant temperatures. The current densities were kept below 3×10^4 A/m^2 in order to suppress flux-flow effects. The characteristics of the transitions in increasing field can be described as a distinct onset, followed by a linear increase, changing to a gradual approach of the normal-state resistance R_n. B_{c2} is defined by extrapolating the linear part to $R=0$ and the transition width ΔB by the intercept between $R=0$ and $R=R_n$. For instance, we obtained for sample II at 5.0 K $B_{c2}=1.51T$ and $\Delta B=0.26T$. It shows that the transitions are relatively broad, indicating the presence of inhomogeneities, probably caused by the splicing procedure. Comparison with Toyota's data[14] on a single crystal with $T_c=7.29$ K supports this conclusion. Our B_{c2} values coincide with his results, but our ΔB values are significantly larger.

The J_c values were determined both from the measured IV curves and from the registration of $J_c(B)$ at a constant-voltage criterion of 2 μV/cm. The characteristic features discussed below do not strongly depend on the criterion as long as it is within a factor of 3 of 2 μV/cm. Typical results for sample II are shown in Fig. 1 for $T=1.8$ and 4.2 K in a plot of $F_p=J_cB$ vs $b=B/B_{c2}$. We have used the B_{c2} values as obtained from the $R(B)$ transitions at the same temperature. Both the characteristic features of 2DCP and of a DCO are clearly observed:[9,12] a smoothly increasing F_p roughly $\propto b^{1/2}$ up to $b\approx0.7$–0.8, a rapid increase (peak effect) up to a well-defined maximum followed by a sharp and almost linear decrease to zero. The size of the peak decreases with increasing temperature.

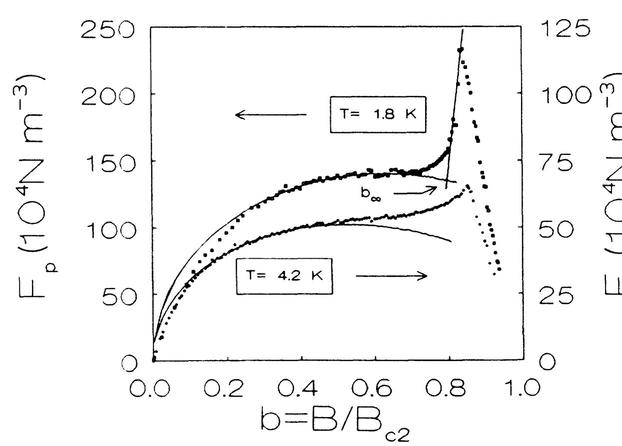

FIG. 1. Pinning force density vs b at $t=0.245$ ($T=1.8$ K) and $t=0.571$ ($T=4.2$ K) for sample II. The solid lines represent the ideal 2DCP behavior. For $T=1.8$ K the construction procedure for $b_{\rm co}$ is shown.

The low-field $F_p(b)$ behavior is typical for 2DCP and this regime provides the information needed to compute R_c and L_c. Using Eqs. (1) and (2) with $r_f = a_0/2$, valid for $b > 0.2$, and $V_c = R_c^2 d$ yields R_c/a_0 and W as a function of b. For $R_c/a_0(b)$ we observed the same dome-shaped behavior in the 2D regime like for amorphous Nb_3Ge,[9] but the maximum values for R_c/a_0, occurring when $b \approx 0.35$, are about 120 for sample I and 240 for sample II, and are therefore much larger than in Nb_3Ge. The large R_c/a_0 values in the 2D regime indicate very large correlated volumes and thus very weak pinning, in good agreement with the scanning-tunneling-microscopy observations by Hess *et al.*[18] which showed a very well-ordered FLL. At the onset field of the peak effect R_c/a_0 has decreased to ≈ 100 for sample II and ≈ 30 for sample I.

For sample II we plotted the results for $W(b)$ as $W/b(1-b)^2$ vs b in Fig. 2 for $T = 1.8$ and 4.2 K. One should realize that the computed data points are also plotted in the high-field regime where the 2DCP theory is not valid. The computed behavior of $W(b)$ is typical[17,19,20] for pinning by defects which both have a δT_c- and a $\delta\kappa$-pinning character, e.g., small, flat precipitates parallel to the layers with different T_c. For such pins W can be expressed as

$$W = [C_1(t) + C_2(t)b]b(1-b)^2 , \quad (6)$$

with[12,19–22]

$$C_1(t) \approx CB_{c2}(t)B_c(t)^2(1+At^2)^2 , \quad (7)$$

and $C_2(t) \propto C_1(t)$. C is a constant depending on the concentration of pins, and $|A|$ a constant of order 1. Figure 2 shows that the behavior according to (6) is observed for $0.2 < b < 0.7$. Below $b = 0.2$ deviations arise because $r_f \approx \xi$ instead of $r_f = a_0/2$ has to be used.[23] The inset of Fig. 2 displays a plot of $(C_1/B_{c2}B_c^2)^{1/2}$ vs t^2 for sample II and it is seen that the result is in accordance with Eq. (7) with $A = 0.5 \pm 0.1$. The ratio C_1/C_2 has an average value -1.3 and is constant within 25%. The above observations indicate that the pinning in $2H$-$NbSe_2$ is indeed caused by small, flat precipitates.

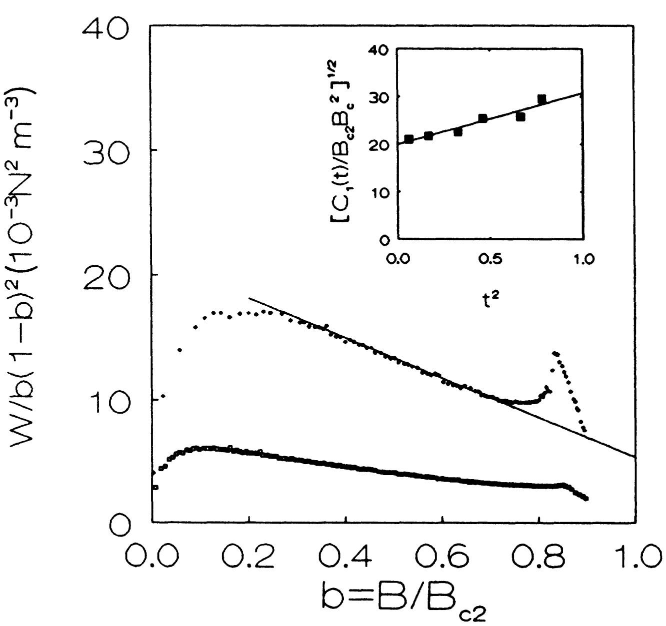

FIG. 2. $W/b(1-b)^2$ vs b at $T = 1.8$ K (upper curve) and $T = 4.2$ K (lower curve) for sample II. The solid line represents the best fit of Eq. (6) to the data in the 2DCP regime. In the inset $(C_1/B_{c2}B_c^2)^{1/2}$ is plotted vs t^2 for the same sample. The solid line represents Eq. (7) with $A = 0.55$.

We finally need to discuss the nature of the peak effect. In the original data (Fig. 1) the onset field for the peak is not sharply defined. Therefore we compute $F_p(b)$ from Eqs. (1), (2), and (6), using fitted values for $C_1(t)$ and $C_2(t)$. The result is given by the solid lines in Fig. 1. Above $b \approx 0.7$, F_p starts to deviate from this ideal 2DCP behavior. We now define the onset field of the peak effect in a manner displayed in Fig. 1 (upper curve). The reason for the gradual transition will be discussed below. The peak effect may indicate the development of plastic deformations in the FLL in the form of edge dislocations.[24] Such defects locally reduce c_{66} so that the FL's can better adjust to the random pin distribution causing an increase of F_p. It has, however, been shown[24] that plastic disorder develops when $R_c \lesssim 17a_0$. Since in our samples R_c/a_0 at the peak onset is considerably larger, see Table I, we conclude that a crossover from two- to three-dimensional FLL disorder is a more likely explanation. Assuming then that the peak effect is caused by a DCO, we computed L_c at b_{co} from Eq. (5) substituting $m_z/m \approx 9$ as deduced from the $H_{c2\parallel}/H_{c2\perp}$ data in Ref. 14. In Table I the values of both b_{co} and the resulting $L_{c,co}/d$ are listed. The error margins, given between parentheses, are related to the uncertainty in b_{co}. The temperature dependence of $L_{c,co}/d$ is similar to that observed for thick amorphous Nb_3Ge films in which a DCO has been unambiguously demonstrated.[9] The agreement with the criterion $L_{c,co} = d/2$ is better for lower temperatures in thicker films. From this we conclude that the peak effect in our $NbSe_2$ samples is caused by a DCO. We note that it is crucial to take into account the anisotropy, for L_c/d would be larger by a factor of 3 if we had used $m/m_z = 1$.

The only question that remains is the different character of the peak observed here and for the DCO in amorphous Nb_3Ge. In the latter case the DCO was characterized by a jump in F_p, whereas in our samples the transi-

TABLE I. L_c/d for several temperatures in both samples at the crossover field b_{co} computed with Eq. (5).

Sample	d (μm)	t	b_{co}	$R_c(b_{co})/a_0$	$L_{c,co}/d$
		0.248	0.82	20	0.63(±0.10)
I	2.0	0.579	0.83	34	1.39(±0.26)
		0.827	···	···	···
		0.245	0.80	100	0.50(±0.08)
		0.408	0.82	101	0.52(±0.09)
II	15.5	0.571	0.76	130	0.80(±0.16)
		0.680	0.75	132	0.87(±0.23)
		0.816	···	···	···
		0.884	···	···	···

tion is rather continuous. In addition, the decay of F_p above the maximum in Fig. 1 is much slower than previously observed.[9] We think this less distinct behavior, like the large ΔT_c and ΔB obtained above, is caused by the inhomogeneities in our single crystals which will smear out all the sharp features and there is no contradiction with our conclusion that the peak effect reveals a DCO.

In summary we have measured the perpendicular field dependence of F_p for several temperatures in very thin single crystals of $2H$-$NbSe_2$. Below b_{co} F_p is well described by the 2DCP theory for an elastically deformed FLL. It is argued that at b_{co} screw dislocations enter the flux-line lattice and destroy the positional correlation along the field direction, thereby changing the disorder in the FLL from 2D to 3D. It is argued that the correlation perpendicular to the field is maintained at b_{co}, and that the peak in F_p is mainly caused by a decrease of L_c. At b_{co} the crossover criterion $L_c = d/2$ is fulfilled with L_c given by Eq. (5), in which for the first time both the electron-mass anisotropy related to the layered structure of our crystal and the dispersion of c_{44} have been taken into account.

We would like to thank Dr. G. A. Wiegers for providing us with single crystals and Professor J. A. Mydosh for his stimulation and interest. This work is part of the research program of the Foundation for the Fundamental Research on Matter.

[1]A. Houghton, R. A. Pelcovits, and A. Subdø, Phys. Rev. B **40**, 6763 (1989).
[2]P. H. Kes, J. Aarts, V. M. Vinokur, and C. J. v.d. Beek, Phys. Rev. Lett. **64**, 1063 (1990).
[3]D. R. Nelson, Phys. Rev. Lett. **60**, 1973 (1988).
[4]E. H. Brandt, Phys. Rev. Lett. **63**, 1106 (1989).
[5]Y. Yeshurun and A. P. Malozemof, Phys. Rev. Lett. **60**, 2202 (1988).
[6]M. V. Feigel'man and V. M. Vinokur, Phys. Rev. B **41**, 8986 (1990).
[7]M. P. A. Fisher, Phys. Rev. Lett. **62**, 1415 (1989).
[8]R. H. Koch, V. Foglietti, W. J. Gallagher, G. Koren, A. Gupta, and M. P. A. Fisher, Phys. Rev. Lett. **63**, 1511 (1989).
[9]R. Wördenweber and P. H. Kes, Physics **135B**, 136 (1985); R. Wördenweber and P. H. Kes, Phys. Rev. B **34**, 494 (1986); P. H. Kes and R. Wördenweber, J. Low Temp. Phys. **67**, 1 (1987); R. Wördenweber and P. H. Kes, Cryogenics **29**, 321 (1989).
[10]E. H. Brandt, J. Low Temp. Phys. **64**, 375 (1986); E. H. Brandt and U. Essman, Phys. Status. Solidi B **144**, 13 (1987); H. R. Kerchner, J. Low Temp. Phys. **50**, 337 (1983).
[11]A. I. Larkin and Yu. N. Ovchinnikov, J. Low Temp. Phys. **34**, 409 (1976).
[12]P. H. Kes and C. C. Tsuei, Phys. Rev. Lett. **47**, 1930 (1981); Phys. Rev. B **28**, 5126 (1983).
[13]B. E. Brown and D. Y. Beernsten, Acta Crystallogr. **18**, 31 (1965).
[14]N. Toyota, H. Nakatsuji, K. Noto, A. Hoshi, N. Kobayashi, Y. Muto, and Y. Onodera, J. Low Temp. Phys. **25**, 485 (1976).
[15]P. de Trey, S. Gygax, and J. P. Jan, J. Low Temp. Phys. **11**, 421 (1973).
[16]E. H. Brandt, J. Low Temp. Phys. **26**, 709 (1977); **26**, 735 (1977); **28**, 263 (1977); **28**, 291 (1977).
[17]P. H. Kes and J. v.d. Berg, in *Studies of High Temperature Superconductors*, edited by A. V. Narlikar (NOVA, New York, 1990), Vol. 5, p. 83.
[18]H. F. Hess, R. B. Robinson, R. C. Dynes, J. M. Valles, Jr., and J. V. Waszczak, Phys. Rev. Lett. **62**, 214 (1989).
[19]P. H. Kes and C. C. Tsuei, J. Non-Cryst. Solids **61-62**, 865 (1984).
[20]R. Wördenweber, A Pruymboom, and P. H. Kes, J. Low Temp. Phys. **70**, 253 (1988); A. Pruymboom, Ph.D. thesis, Leiden University, 1988 (unpublished).
[21]C. S. Pande, Appl. Phys. Lett. **28**, 462 (1976).
[22]E. J. Kramer, Philos. Mag. **33**, 331 (1976).
[23]E. H. Brandt, Phys. Rev. Lett. **57**, 1347 (1986).
[24]R. Wördenweber, P. H. Kes, and C. C. Tsuei, Phys. Rev. B **33**, 3172 (1986).